

Neutron-star mergers in scalar-tensor theories of gravity

Enrico Barausse,^{1,2} Carlos Palenzuela,³ Marcelo Ponce,² and Luis Lehner^{4,5}

¹*Institut d'Astrophysique de Paris/CNRS, 98bis boulevard Arago, 75014 Paris, France*

²*Department of Physics, University of Guelph, Guelph, Ontario N1G 2W1, Canada*

³*Canadian Institute for Theoretical Astrophysics, Toronto, Ontario M5S 3H8, Canada*

⁴*Perimeter Institute for Theoretical Physics, Waterloo, Ontario N2L 2Y5, Canada*

⁵*CIFAR, Cosmology & Gravity Program, Canada*

(Dated: December 21, 2012)

Scalar-tensor theories of gravity are among the most natural phenomenological alternatives to General Relativity, because the gravitational interaction is mediated by a scalar degree of freedom, besides the gravitons. In regions of the parameter space of these theories where constraints from both solar system experiments and binary-pulsar observations are satisfied, we show that binaries of neutron stars present marked differences from General Relativity in both the late-inspiral and merger phases. These strong-field effects are difficult to reproduce in General Relativity, even with an exotic equation of state. We comment on the relevance of our results for the upcoming Advanced LIGO/Virgo detectors.

PACS numbers:

General Relativity (GR) has passed stringent tests in the solar system [1] and in binary pulsars [2]. However, these tests involve weak gravitational fields and/or velocities $v \ll c$, so the theory remains essentially untested in the strong-field $v \sim c$ regime, where high-energy corrections may appear. Because strong-field regimes are naturally defined by systems containing black holes (BHs) and/or neutron stars (NS's), the final stages in the evolution of compact-object binaries provide the best opportunities to explore gravitation at extreme conditions [3].

Astrophysical evidence for the existence of BHs, albeit convincing, is still not accurate enough to discriminate between GR and possible alternatives. Insight in the nature of BHs, and on whether they are the objects predicted by GR, may be provided by future electromagnetic probes (see e.g. [4–8]). Gravitational waves (GWs) provide cleaner prospects for this task, because their generation and propagation is largely insensitive to the complicated astrophysical environment in which BHs live. GW-based tests of BHs and gravity theories will be possible with the upcoming Advanced LIGO/Virgo detectors [9–12] and to an exquisite degree with future detectors such as LISA [9, 11, 13–19].

Tests of gravity theories with electromagnetic and GW probes will also be (and to some extent, already are) possible with NS's. Interestingly, these objects are generally more sensitive than BHs to the presence of extra degrees of freedom in the gravity theory (see Refs. [20–24] for some representative examples). Such behavior has allowed for placing constraints on gravitational theories using observations of isolated NS's [20–24] or widely separated binary pulsars evolving under the effect of GW emission [25–29]. More stringent constraints may be possible with the expected Advanced LIGO/Virgo detection of binary NS's and BH-NS binaries [30, 31].

In this Letter, we consider binary NS systems and fo-

cus on strong-field effects during the late inspiral/plunge until merger of the two stars (after merger effects are deferred to a future work). Although the merger is only marginally detectable with Advanced LIGO/VIRGO in GR, we will show and discuss here that if one considers modifications to the gravity theory: (i) strong discrepancies arise which are not captured by the simpler weak-field analyses; (ii) these effects cannot be reproduced within GR, even with an exotic equation of state; (iii) observable features will be detectable with Advanced LIGO/VIRGO, even in the late inspiral/plunge and merger, unlike what happens if the gravity theory is unmodified GR; (iv) these features may even have astrophysical implications in possible models for energetic electromagnetic events.

For the gravity theory, we focus here on scalar-tensor (ST) theories [32–34], in which the gravitational field is described by the usual tensor degrees of freedom and a non-minimally coupled scalar field. These viable alternative theories of gravity can also be regarded as natural ones, because they include an extra scalar degree of freedom in the gravitational field, as expected based on string theory. As well, many phenomenological gravity theories can be either shown to be exactly equivalent to a ST theory (c.f. for instance $f(R)$ gravity [32, 35]) or to contain a gravitational scalar besides other degrees of freedom. ST theories are also among the most strongly constrained alternatives to GR, since their history dates back to the 50-60's with Jordan [36], Fierz [37], Brans and Dicke [38] (whose theory is a particular ST theory). Bounds have been placed on these theories with solar system experiments [1], isolated NS [20, 39] and binary pulsars [25–28]. Stricter constraints may be obtained by detecting the GW signal from a gravitational collapse [40] or from vibrating NS's [41]. The remaining viable parameter space of ST theories is nevertheless still sizable, and these the-

ories are therefore a rather natural choice to investigate strong-field deviations from GR. In fact, we will show that for viable ST theories and NS binaries, strong field effects are possible that are qualitatively different from GR, and which are related to the “spontaneous scalarization” of isolated NS’s in ST theories, first discovered in Ref. [20].

Methodology: We consider a generic ST theory

$$S = \frac{1}{2\kappa} \int d^4x \sqrt{-g} \left[\phi R - \frac{\omega(\phi)}{\phi} \partial_\mu \phi \partial^\mu \phi - V(\phi) \right] + S_M(g_{\mu\nu}, \psi), \quad (1)$$

where $\kappa = 8\pi G$ (we are setting $c = 1$ throughout this Letter), R and g are respectively the Ricci scalar and determinant of the metric, ϕ is the extra scalar field describing the gravitational field, $V(\phi)$ is a generic potential, and ψ collectively describes the matter degrees of freedom. Jordan-Fierz-Brans-Dicke theory corresponds to the particular case $\omega = \text{const}$, while for $\omega = 0$ ($\omega = -3/2$) and a suitable potential, the theory reduces to metric $f(R)$ (Palatini $f(R)$) gravity. Theories with $\omega(\phi) = -3/2 - \kappa/(4\beta \log \phi)$ are equivalent (as long as $\phi > 1$) to those studied in Ref. [20], which give significant deviations from GR for spherical NS’s (“spontaneous scalarization”).

A somewhat simpler form for ST theories is obtained by re-expressing the (“Jordan-frame”) action (1) into the so-called “Einstein-frame” action through a conformal transformation $g_{\mu\nu}^E = \phi g_{\mu\nu}$, which yields

$$S = \int d^4x \sqrt{-g^E} \left[\frac{R^E}{2\kappa} - \frac{1}{2} g_E^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi - \tilde{V}(\varphi) \right] + S_M(g_{\mu\nu}^E/\phi(\varphi), \psi), \quad (2)$$

where $\tilde{V}(\varphi) = V(\phi)/(2\kappa\phi^2)$, and the scalar field φ is related to ϕ by $d \log \phi / d\varphi = \{2\kappa/[3 + 2\omega(\phi)]\}^{1/2}$. Imposing $\varphi = 0$ for $\phi = 1$, this can be integrated to give

$$\phi = \exp \left(\sqrt{\frac{2\kappa}{3 + 2\omega}} \varphi \right), \quad \phi = \exp(-\beta \varphi^2), \quad (3)$$

respectively for Jordan-Fierz-Brans-Dicke and for the theories considered in Ref. [20]. (Note that our φ is related to the scalar field φ_{DEF} used by Ref. [20] via $\varphi = \varphi_{\text{DEF}}/\sqrt{4\pi G}$.)

In the Einstein frame the field equations are

$$G_{\mu\nu}^E = \kappa (T_{\mu\nu}^\varphi + T_{\mu\nu}^E), \quad (4)$$

$$\square^E \varphi - \tilde{V}'(\varphi) = \frac{1}{2} \sqrt{\frac{2\kappa}{3 + 2\omega(\phi)}} T_E, \quad (5)$$

$$\nabla_\mu^E T_E^{\mu\nu} = -\frac{1}{2} T_E \sqrt{\frac{2\kappa}{3 + 2\omega(\phi)}} g_E^{\mu\nu} \partial_\mu \varphi, \quad (6)$$

where

$$T_E^{\mu\nu} = \frac{2}{\sqrt{-g^E}} \frac{\delta S_M}{\delta g_{\mu\nu}^E} \quad \text{and} \quad (7)$$

$$T_{\mu\nu}^\varphi = \partial_\mu \varphi \partial_\nu \varphi - g_{\mu\nu}^E \left[\frac{1}{2} g_E^{\alpha\beta} \partial_\alpha \varphi \partial_\beta \varphi + \tilde{V}(\varphi) \right] \quad (8)$$

are the scalar-field and matter stress-energy tensors in the Einstein frame, and $T_E \equiv T_E^{\mu\nu} g_{\mu\nu}^E$. The indices of the matter stress-energy tensor are raised/lowered with the Einstein-frame metric g_E , and the relation to the Jordan-frame stress-energy tensor is given by $T_E^{\mu\nu} = T^{\mu\nu} \phi^{-3}$, $T_{\mu\nu}^E = T_{\mu\nu} \phi^{-1}$. Also, the relations between the matter variables in the two frames are $u^\mu = \sqrt{\phi} u_E^\mu$ (from the normalization condition $g_{\mu\nu}^E u_E^\mu u_E^\nu = -1$); $\rho = \phi^2 \rho_E$ (from $\rho_E = u_E^\mu u_E^\mu T_{\mu\nu}^E$) and $p = \phi^2 p_E$ (from $T = \phi^2 T_E$). Last, in order to have the same equation of state in both frames, one must have $\rho_0 = \phi^2 \rho_0^E$. This definition, together with the current conservation in the Jordan frame ($\nabla_\mu j^\mu = 0$ with $j^\mu = \rho_0 u^\mu$), gives

$$\nabla_\mu^E j_E^\mu = -\frac{1}{2\phi} j_E^\mu \partial_\mu \phi = -\frac{1}{2} j_E^\mu \partial_\mu \varphi \sqrt{\frac{2\kappa}{3 + 2\omega(\phi)}}, \quad (9)$$

with $j_E^\mu = \rho_0^E u_E^\mu$. Therefore, solving the system given by Eqs. (4), (5), (6) and (9) and transforming back to the original Jordan frame provides a solution to the original problem. This is the approach we follow in this work.

Physical Set-up: We model the NS’s with a perfect fluid coupled to the full field equations (4-6, 9) to accurately represent the strong gravitational effects during the evolution of a binary system. Our numerical techniques for solving these coupled equations have been thoroughly described and tested previously [42–47]. The initial data are evolved in a cubical computational domain defined by $x^i \in [-350, 350]$ km, and we employ an adaptive mesh refinement that tracks the two compact objects with cubes slightly larger than the radius of star and resolution $\Delta x = 0.5$ km.

We consider an unequal-mass binary system, initially on a quasi-circular orbit with separation of 60 km and angular velocity $\Omega = 1295$ rad/s, constructed with LORENE [48]. The stars are described by a polytropic equation of state ($p/c^2 = K \rho_0^\Gamma$) with $\Gamma = 2$ and $K = 123 G^3 M_\odot^2 / c^6$. We adopt a mass ratio of $q \equiv 0.937$, possible for progenitors of gamma-ray bursts [49], and choose individual baryon masses of $\{1.78, 1.90\} M_\odot$, corresponding to gravitational masses $\{1.58, 1.67\} M_\odot$.

For the gravity theory, we consider $V(\phi) = 0$ and $\omega(\phi) = -3/2 - \kappa/(4\beta \log \phi)$. As mentioned, these theories are equivalent to those studied in Ref. [20] as long as $\phi > 1$ (or equivalently, $\varphi > 0$). Because the coupling ω diverges when $\phi = 1$, we initially set $\phi > 1$, and check that this condition is satisfied throughout our evolutions. This guarantees the exact equivalence between our setup and that of Ref. [20]. Besides the constant β ,

the gravity theory is also characterized by the asymptotic value φ_0 of the scalar field far from the source [20]. Binary pulsars measurements require $\beta/(4\pi G) \gtrsim -4.5$ [26], while the Cassini experiment constrains $\varphi_0 < \varphi_{\text{Cassini}} \equiv 2(G\pi)^{1/2}/[|\beta|(3+2\omega_0)^{1/2}] \approx 1.26 \times 10^{-2} G^{1/2}/|\beta|$ (with $\omega_0 = 4 \times 10^4$ [1, 26]). Moreover, from $\beta/(4\pi G) \sim -4$ to $\beta/(4\pi G) = -4.5$, the allowed value for φ_0 decreases monotonically from φ_{Cassini} to 0, again due to constraints from binary pulsars [26]. In our simulations, we tried different allowed values of φ_0 , and the results do not change significantly when β is fixed.

GW extraction and backreaction: The response of a GW detector far from the source is encoded in the curvature scalars in the physical (Jordan) frame [50]. These are readily obtained from the Einstein frame components as $\psi_4 = -R_{\ell\bar{m}\ell\bar{m}} = \phi\psi_4^E$, $\psi_3 = -R_{\ell k\ell\bar{m}}/2 = \phi\psi_3^E + \dots$, $\psi_2 = -R_{\ell k\ell k}/6 = \phi\psi_2^E + \dots$ and $\phi_{22} = -R_{\ell m\ell\bar{m}} = \phi(\phi_{22}^E - l^\nu l^\mu \nabla_\nu \nabla_\mu \log \phi/2 + \dots)$ (with \dots denoting sub-leading terms in the distance to the detector and l, m being components of a null tetrad adapted to outgoing wavefronts). Because far from the source one expects $\varphi = \varphi_0 + \varphi_1/r + \mathcal{O}(1/r^2)$, with $\varphi_0 = \text{const}$ and φ_1 a function of x^μ , and because of the peeling property in the Einstein frame, it is clear that ψ_2 and ψ_3 decay faster than $1/r$ and do not produce observable effects on a GW detector at infinity. However, using the fact that $\log \phi = -\beta\varphi^2 = -\beta(\varphi_0^2 + 2\varphi_0\varphi_1/r) + \mathcal{O}(1/r^2)$, one easily obtains $\phi_{22} \sim \beta\varphi_0\partial_t^2\varphi_1/r$. As a result, the radiative degrees of freedom (decaying as $1/r$ and therefore observable by GW detectors) are ψ_4 (carrying tensorial degrees of freedom) and $\phi_{22} \sim \beta\varphi_0\partial_t^2\varphi_1/r$ (carrying a purely transverse, radiative scalar mode [50]).

It should be noted, however, that for $\varphi_0 \rightarrow 0$ the $1/r$ radiative component of ϕ_{22} vanishes. This is physically relevant because, as mentioned, φ_0 is constrained to small values and means that for viable ST theories the purely transverse scalar modes couple rather weakly to GW detectors. This makes their direct detection problematic; which at first sight seems odd as eqs. (6) and (5) imply that the scalar field carries energy to infinity. (This is because the energy flux is essentially given by the integral of T^{tr} on a 2-sphere at infinity, and $T^{tr} \sim \partial_t^2\varphi_1/r^2$). It is easy, however, to get convinced directly at the level of the action (2) that these fluxes are not observable with GW detectors in the limit $\varphi_0 \rightarrow 0$. The detection of GWs is based on free-falling test masses, so to analyze the detector's response one needs to look at the Jordan frame metric $g_{\mu\nu}^E/\phi(\varphi)$, to which the matter fields ψ couple [cf. eq. (2)]. Far from the source, in suitable coordinates one has $g_{\mu\nu}^E \approx \eta_{\mu\nu} + h_{\mu\nu}$ and $\varphi \approx \varphi_0 + \delta\varphi$, where $h_{\mu\nu}$ and $\delta\varphi$ are small perturbations. If $\varphi_0 = 0$, we have $\phi = \exp(-\beta\varphi^2) \approx 1 - \beta\delta\varphi^2$, and therefore $g_{\mu\nu}^E/\phi(\varphi) \approx \eta_{\mu\nu} + h_{\mu\nu}$ at linear order. This means that the motion of the detector's test masses is only sensitive to the tensor waves $h_{\mu\nu}$ in the limit $\varphi_0 \rightarrow 0$.

Although weakly coupled to GW detectors, the scalar

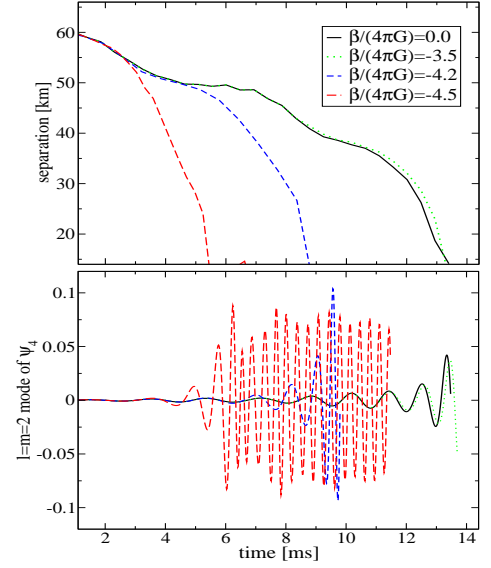


FIG. 1: The binary's separation and the dominant mode of the ψ_4 scalar (encoding the effect of GWs) for different values of β .

fluxes exert a significant backreaction on the source, because they appear at 1.5PN order, while the quadrupolar tensor waves of GR appear at 2.5PN. More precisely, if we consider a quasicircular binary of stars with masses m_1 and m_2 , and scalar charges α_1 and α_2 [with $\alpha_i = \sqrt{4\pi G}\varphi_1^i/m_i$, where the mass scale φ_1 is defined, as above, by $\varphi = \varphi_0 + \varphi_1/r + \mathcal{O}(1/r^2)$], the dipole scalar emission is [25, 27, 39]

$$\dot{E}_{\text{dipole}} = \frac{G}{3c^3} \left(\frac{G_{\text{eff}} m_1 m_2}{r^2} \right)^2 (\alpha_1 - \alpha_2)^2. \quad (10)$$

Here, $G_{\text{eff}} = G(1 + \alpha_1\alpha_2)$ is the effective gravitational constant appearing in the Newtonian interaction between the two stars, i.e. the gravitational force gets modified due to the exchange of scalar gravitons and becomes [39]

$$F = \frac{G_{\text{eff}} m_1 m_2}{r^2}. \quad (11)$$

The quadrupole tensor emission is instead [27, 39]

$$\dot{E}_{\text{quadrupole}} = \frac{32G}{5c^3} \left(\frac{G_{\text{eff}} m_1 m_2}{r^2} \right)^2 \left(\frac{v}{c} \right)^2, \quad (12)$$

where $v = [G_{\text{eff}}(m_1 + m_2)/r]^{1/2}$ is the relative velocity of the two stars. Therefore, the dipole scalar fluxes, although directly undetectable, are produced abundantly during the binary's inspiral if the charges α_1 and α_2 are different, and dominate over the tensor quadrupole fluxes, which are suppressed by $(v/c)^2$ relative to them.

Numerical evolutions and comparison to GR: Our simulations confirm the qualitative features described above, but also highlight a more intricate phenomenology. During the inspiral phase, scalar fluxes are emitted if the

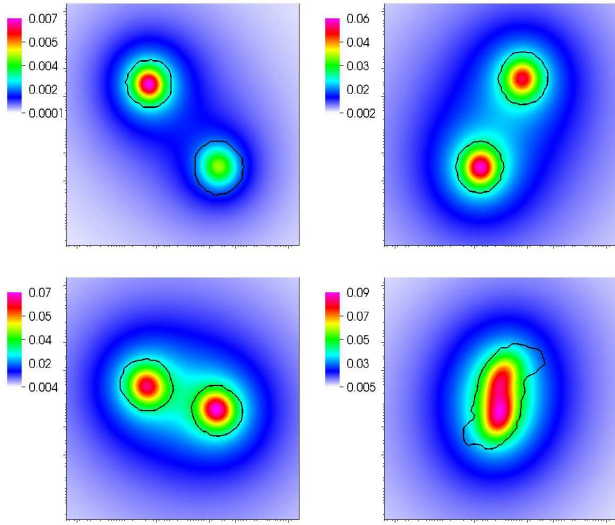


FIG. 2: The scalar field (color code) and the NS surfaces (solid black line) at $t = \{1.8, 3.1, 4.0, 5.3\}$ ms for $\beta/(4\pi G) = -4.5$.

scalar charges of the two stars differ. These fluxes are difficult to detect directly, but clearly affect the orbital evolution, and their effect is thus encoded in the curvature scalar ψ_4 , i.e. in the tensor GWs (cf. bottom panel of Fig. 1). Specifically, in ST-theories with $\beta/(4\pi G) \lesssim -4.2$, NS binaries merge at lower frequency than in GR, e.g. in Fig. 1 the plunge starts already when the stars’ centers are ~ 52 km apart, corresponding to an angular velocity $\Omega \sim 1850$ rad/s (i.e. a GW frequency $f \sim \Omega/\pi \sim 586$ Hz, within Advanced LIGO/Virgo’s sensitivity bands), and results in the formation of rotating bar (whose long-lived GW signal can be seen in the lower panel). Remarkably, plunges starting so early cannot be obtained in GR, because even with exotic equations of state, NS radii are constrained to $R \lesssim 14$ km [51], so the interaction between the two stars cannot trigger a plunge until a separation $\sim 2R \lesssim 28$ km.

The cause of these earlier merger is not simply the backreaction of the scalar fluxes (10) (absent in GR). In fact, we observe that even if we attempt to maximize the dipole emission (10) by setting up a star with the maximum allowed scalar charge for a given ST theory ($\alpha_1 = \alpha_{\text{max}}$) and one with zero scalar charge ($\alpha_2 \approx 0$), the scalar field grows rapidly inside the initially non-scalarized star, which quickly develops a charge $\alpha_2 \approx \alpha_1$ when the binary becomes sufficiently close (cf. Fig. 2). This shuts off the dipole flux (10), but enhances the Newtonian force pulling the stars together, eq. (11). Therefore the more rapid mergers in ST theories are caused by the combination of dissipative [eq. (10)] and conservative [eq. (11)] effects. Indeed, as a qualitative test, we have performed a direct integration of the PN equations of motion of GR with the gravitational constant G replaced by $G_{\text{eff}} = G(1 + \alpha_1\alpha_2)$ [so as to mimic eq. (11),

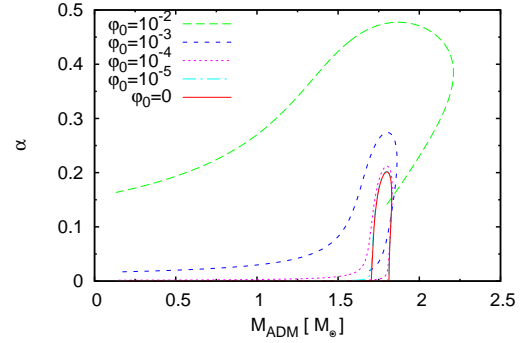


FIG. 3: Effect of an external scalar field φ_0 , for $\beta/(4\pi G) = -4.5$.

with $\alpha_1, \alpha_2 \sim 0.2 - 0.4$ set to values compatible with our simulations], and this exercise confirms that the enhanced gravitational pull is enough to produce much earlier mergers.

The observed growth of the scalar field and charge of initially non-scalarized stars that get close to scalarized ones can be understood in simple terms. (This phenomenon is known as “induced scalarization” [20, 39], and has also been observed for boson stars in ST theory [52].) The scalar field of scalarized stars extends well beyond the radius of the baryonic matter [20]. When the non-scalarized star enters this scalar-field “halo” of the scalarized star, it grows a significant scalar charge. This can be seen by studying isolated NS’s in ST theory [20], and imposing a non-zero asymptotic value φ_0 for the scalar field, in order to mimic the effect of the “external” scalar field produced by the halo of the other (scalarized) star. The effect of φ_0 is shown in Fig. 3, where we used a static, spherically symmetric code to calculate the scalar charge of NS’s as a function of the ADM mass, for a ST theory with $\beta/(4\pi) = -4.5$. As can be seen, even modest values of φ_0 induce significant scalar charges. This is similar, energetically, to the magnetization of a ferromagnetic material immersed in a sufficiently strong magnetic field [53]. Here, the external scalar field makes the configuration with non-zero charge energetically preferred over the initial non-charged one.

Finally, the total ADM mass of the systems we consider is $3.35M_{\odot}$, slightly larger than the masses of observed binary pulsars ($\sim 2.8 - 3M_{\odot}$), but still plausible. This value was mainly chosen to maximize the deviations from the pure-GR case, but the effects we describe in this Letter are expected to arise, although in a weaker fashion, also for binaries with lower masses, which we plan to study in a follow-up work. Also, NS binaries are among the most likely progenitors for short gamma-ray bursts (GRBs), so irrespective of the properties of the currently observed binaries, it makes sense to investigate if modifications of the gravity theory can affect their physics. Our findings show that in ST theories the GW signal accompanying GRBs may be weaker (because part of the energy is car-

ried away in scalar waves) and different from GR, which would make it tricky to detect, even in the presence of an electromagnetic counterpart. This may have important implications for coincident searches of GW and electromagnetic signals from GRBs. More generally, an interesting consequence of our results is that for sufficiently massive binaries the late-time orbiting and merger phases as well as the after-merger object and its dynamics can be significantly different from GR. This can strongly affect energetic events possibly associated with such mergers.

Acknowledgments: We are deeply indebted to Gilles Esposito-Farese for providing insightful suggestions on different physical effects, observables and on the spontaneous scalarization induced on a star by another in ST theory. Also, we thank our collaborators M. Anderson, E. Hirschmann, S.L. Liebling and D. Neilsen with whom we have developed the basic computational infrastructure employed in this work. E. B. acknowledges support from a CITA National Fellowship while at the University of Guelph, and partial support from a Marie Curie Career Integration Grant (PCIG11-GA-2012-321608) while at the Institut d'Astrophysique de Paris. C. P. acknowledges support from the Jeffrey L. Bishop Fellowship. L. L. acknowledges support from NSERC through a Discovery grant. All authors also acknowledge hospitality from the Kavli Institute for Theoretical Physics (UCSB), where part of this work was carried out. This work was supported in part by the National Science Foundation under grant No. NSF PHY11-25915. Computations were performed on Scinet. Research at Perimeter Institute is supported through Industry Canada and by the Province of Ontario through the Ministry of Research & Innovation.

-
- [1] C. M. Will, Living Rev. Rel. **4**, 4 (2001) [gr-qc/0103036].
 - [2] J. H. Taylor and J. M. Weisberg, J. M. Astrophys. J. **253**, 908 (1982); T. Damour and J. H. Taylor, Phys. Rev. D **45**, 1840 (1992).
 - [3] J. Abadie *et al.* [LIGO Scientific and Virgo Collaborations], Class. Quant. Grav. **27**, 173001 (2010).
 - [4] C. M. Will, Astrophys. J. Lett. **674**, L25 (2008);
 - [5] S. Doeleman, E. Agol, D. Backer, F. Baganoff, G. C. Bower, A. Broderick, A. Fabian and V. Fish *et al.*, arXiv:0906.3899 [astro-ph.CO].
 - [6] C. Bambi, E. Barausse, Astrophys. J. **731**, 121 (2011); Phys. Rev. D **84**, 084034 (2011);
 - [7] T. Johannsen and D. Psaltis, Astrophys. J. **716**, 187 (2010); **718**, 446 (2010); Astrophys. J. **726**, 11 (2011); Astrophys. J. **745**, 1 (2012);
 - [8] C. Bambi, K. Freese, Phys. Rev. D **79**, 043002 (2009); C. Bambi and N. Yoshida, Class. Quant. Grav. **27**, 205006 (2010).
 - [9] N. Yunes and F. Pretorius, Phys. Rev. D **80**, 122003 (2009); N. Cornish, L. Sampson, N. Yunes and F. Pretorius, Phys. Rev. D **84**, 062003 (2011)
 - [10] C. K. Mishra, K. G. Arun, B. R. Iyer and B. S. Sathyaprakash, Phys. Rev. D **82**, 064010 (2010)
 - [11] S. Mirshekari, N. Yunes and C. M. Will, Phys. Rev. D **85** (2012) 024041 [arXiv:1110.2720 [gr-qc]].
 - [12] T. G. F. Li, W. Del Pozzo, S. Vitale, C. Van Den Broeck, M. Agathos, J. Veitch, K. Grover and T. Sidery *et al.*, Phys. Rev. D **85**, 082003 (2012)
 - [13] C. M. Will and N. Yunes, Class. Quant. Grav. **21**, 4367 (2004); E. Berti, A. Buonanno and C. M. Will, Phys. Rev. D **71** (2005) 084025
 - [14] V. Cardoso, S. Chakrabarti, P. Pani, E. Berti and L. Gualtieri, Phys. Rev. Lett. **107**, 241101 (2011); N. Yunes, P. Pani and V. Cardoso, Phys. Rev. D **85**, 102003 (2012).
 - [15] F. D. Ryan, Phys. Rev. D **52**, 5707 (1995); **56**, 1845 (1997); **56**, 7732 (1997); K. Glampedakis and S. Babak, Class. Quant. Grav. **23**, 4167 (2006); L. Barack and C. Cutler, Phys. Rev. D **75**, 042003 (2007)
 - [16] T. A. Apostolatos, G. Lukes-Gerakopoulos and G. Contopoulos, Phys. Rev. Lett. **103**, 111101 (2009);
 - [17] M. Kesden, J. Gair and M. Kamionkowski, Phys. Rev. D **71**, 044015 (2005);
 - [18] C. F. Sopuerta and N. Yunes, Phys. Rev. D **80**, 064006 (2009); P. Canizares, J. R. Gair and C. F. Sopuerta, Phys. Rev. D **86**, 044010 (2012); K. Yagi, N. Yunes and T. Tanaka, arXiv:1208.5102 [gr-qc].
 - [19] E. Berti, V. Cardoso and C. M. Will, Phys. Rev. D **73**, 064030 (2006)
 - [20] T. Damour and G. Esposito-Farese, Phys. Rev. Lett. **70**, 2220 (1993)
 - [21] P. Pani, V. Cardoso, E. Berti, J. Read and M. Salgado, Phys. Rev. D **83**, 081501 (2011)
 - [22] P. Pani, E. Berti, V. Cardoso and J. Read, Phys. Rev. D **84**, 104035 (2011).
 - [23] C. Eling, T. Jacobson and M. Coleman Miller, Phys. Rev. D **76**, 042003 (2007) [Erratum-ibid. D **80**, 129906 (2009)]
 - [24] E. Barausse, T. P. Sotiriou and J. C. Miller, Class. Quant. Grav. **25**, 062001 (2008); Class. Quant. Grav. **25**, 105008 (2008)
 - [25] D. M. Eardley, Astrophys. J. Lett. **196**, L59 (1975)
 - [26] T. Damour and G. Esposito-Farese, Phys. Rev. D **58**, 042001 (1998); P. C. C. Freire, N. Wex, G. Esposito-Farese, J. P. W. Verbiest, M. Bailes, B. A. Jacoby, M. Kramer and I. H. Stairs *et al.*, Mon. Not. Roy. Astron. Soc. **423**, 3328 (2012)
 - [27] C. M. Will and H. W. Zaglauer, Astrophys. J. **346**, 366 (1989)
 - [28] J. Alsing, E. Berti, C. M. Will and H. Zaglauer, Phys. Rev. D **85**, 064041 (2012)
 - [29] B. Z. Foster, Phys. Rev. D **73**, 104012 (2006) [Erratum-ibid. D **75**, 129904 (2007)]
 - [30] E. Berti, L. Gualtieri, M. Horbatsch and J. Alsing, Phys. Rev. D **85**, 122005 (2012).
 - [31] K. G. Arun, Class. Quant. Grav. **29**, 075011 (2012).
 - [32] R. Wagoner, Phys. Rev. D **1**, 3209 (1970).
 - [33] P. G. Bergmann, Int. J. Theor. Phys. **1**, 25 (1968).
 - [34] K. Nordtvedt, Astrophys. J. **161**, 1059 (1970)
 - [35] T. P. Sotiriou, Class. Quant. Grav. **23**, 5117 (2006).
 - [36] P. Jordan, Z. Phys. **157**, 112 (1959)
 - [37] M. Fierz, Helv. Phys. Acta **29**, 128 (1956)
 - [38] C. Brans and R. H. Dicke, Phys. Rev. **124**, 925 (1961)
 - [39] T. Damour and G. Esposito-Farese, Class. Quant. Grav. **9**, 2093 (1992).
 - [40] J. Novak, Phys. Rev. D **57**, 4789 (1998).
 - [41] H. Sotani and K. D. Kokkotas, Phys. Rev. D **70**, 084026

- (2004)
- [42] M. Anderson, E. Hirschmann, S. L. Liebling and D. Neilsen, *Class. Quant. Grav.* **23**, 6503 (2006) [gr-qc/0605102].
 - [43] C. Palenzuela, I. Olabarrieta, L. Lehner and S. L. Liebling, *Phys. Rev. D* **75**, 064005 (2007).
 - [44] S. L. Liebling, *Phys. Rev. D* **66**, 041703 (2002).
 - [45] M. Anderson, E. W. Hirschmann, L. Lehner, S. L. Liebling, P. M. Motl, D. Neilsen, C. Palenzuela and J. E. Tohline, *Phys. Rev. D* **77**, 024006 (2008).
 - [46] M. Anderson, E. W. Hirschmann, L. Lehner, S. L. Liebling, P. M. Motl, D. Neilsen, C. Palenzuela and J. E. Tohline, *Phys. Rev. Lett.* **100**, 191101 (2008)
 - [47] <http://www.had.liu.edu/>.
 - [48] <http://www.lorene.obspm.fr>.
 - [49] K. Belczynski, R. W. O’Shaughnessy, V. Kalogera, F. Rasio, R. Taam and T. Bulik, *Astrophys. J. Lett.* **680** L129 (2008)
 - [50] D. M. Eardley, D. L. Lee and A. P. Lightman, *Phys. Rev. D* **8**, 3308 (1973)
 - [51] A. W. Steiner, J. M. Lattimer and E. F. Brown, *Astrophys. J.* **722**, 33 (2010)
 - [52] M. Ruiz, J. C. Degollado, M. Alcubierre, D. Nunez and M. Salgado, *Phys. Rev. D* **86**, 104044 (2012)
 - [53] G. Esposito-Farese, *AIP Conf. Proc.* **736**, 35 (2004) [gr-qc/0409081].